

Impact of FCNC top quark interactions on $BR(t \rightarrow bW)$

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We study the effect that FCNC interactions of the top quark will have on the branching ratio of charged decays of the top quark. We have performed an integrated analysis using Tevatron and B-factories data and with just the further assumption that the CKM matrix is unitary we can obtain very restrictive bounds on the strong and electroweak FCNC branching ratios $Br(t \rightarrow qX) < 4.0 \times 10^{-4}$, where X is any vector boson and a sum in $q = u, c$ is implied.

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I. INTRODUCTION

With the large statistics of top quark production expected at the LHC - both in the $t\bar{t}$ and single top channels - precision studies of the properties of this particle will become possible. The detailed study of the branching ratios of the decays of the top is of great interest, since those observables may change by several orders of magnitude, depending on which model we are considering: the Standard Model (SM) or some of its extensions, such as supersymmetry (SUSY) or the two-higgs doublet models (2HDM). For instance, the branching ratio for the flavour-changing neutral current (FCNC) decay $t \rightarrow cg$ is expected to be of the order of 10^{-12} in the SM, whereas in some realizations of SUSY models it may reach as much as 10^{-4} [1, 2]. Thus, a measurement of such branching ratios may be an excellent way of discovering new physics beyond that of the SM.

Recently, in a series of works [3, 4], we have studied the FCNC interactions of the top quark, using the effective operator formalism of Buchmüller and Wyler [6]. That formalism allows us to parameterize whatever new physics may exist beyond the SM in a model-independent manner, by using operators of dimension larger than four which obey the gauge symmetries of the SM. That is, indeed, the greatest advantage of adopting this formalism: gauge invariance is automatically guaranteed. In refs. [3], only operators which induced FCNC in the strong interactions of the top quark were considered. They manifested themselves in FCNC decays of the form $t \rightarrow ug$ and $t \rightarrow cg$, and their impact on processes of top production at the LHC was studied in detail. It was shown that these FCNC interactions might have considerable impact on processes of single top production and associated production of a top quark plus a photon, a Z boson or a Higgs scalar. In ref. [4] we analysed in detail other sources of FCNC processes, namely in the electroweak sector, with new possible decays such as $t \rightarrow u\gamma$ or $t \rightarrow cZ$. We showed that the top electroweak FCNC interactions could contribute as much as the strong ones for processes such as $t + Z$ production at the LHC. We also discussed the possibility of experimentally distinguishing both types of FCNC contributions. An analysis of the effect of all of these FCNC interactions in the production of $t + \text{jet}$ at the LHC has recently been concluded [5].

In this work we wish to study the contributions that these FCNC interactions may have on a more basic top observable: the branching ratio of $t \rightarrow bW$, which is expected, in the SM, to be very close to 1. An obvious way in which the top FCNC interactions will affect $BR(t \rightarrow bW)$ is by changing the total top width: since $BR(t \rightarrow bW) = \Gamma(t \rightarrow bW)/\Gamma_t(\text{total})$, with anomalous FCNC decays allowed the

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value of the total top width will surely increase, thus, in principle, decreasing the value of $BR(t \rightarrow bW)$. However, we will show in this paper the existence of a more curious effect, namely, that gauge invariance implies that the same operators which generate FCNC top decays also give contributions to charged top decays of the form $t \rightarrow dW$ and $t \rightarrow sW$.

This paper is organized as follows: in section II we will briefly review the effective operator formalism and the criteria behind the selection of those operators for FCNC purposes. We will also review the results of [3, 4] concerning the FCNC decay widths of the top quark. In section III we will present the contributions of the FCNC operators to the charged top decays, and the calculation of the branching ratio for $t \rightarrow bW$. In section IV we will show predictions for the FCNC cross sections and branching ratios in terms of $BR(t \rightarrow bW)$ and discuss the importance that precision measurements of this observable might have on the cross section of processes such as $t + Z$ and $t + jet$ production at the LHC. An overview of our conclusions is presented in section V.

II. EFFECTIVE OPERATORS CONTRIBUTING TO TOP FCNC INTERACTIONS

The effective operator formalism [6] considers that, whatever new physics may exist beyond the SM, it will reflect itself at low energies as operators of dimension larger than four, built with the same fields of the SM and respecting the gauge symmetries $SU(3)_C \times SU(2)_W \times U(1)_Y$. The SM is thus seen as the low-energy limit of a more general theory, the effects of which are described in a model-independent manner by operators of dimension 5, 6, etc. Hence, the lagrangian of the theory becomes

$$\mathcal{L} = \mathcal{L}^{SM} + \frac{1}{\Lambda} \mathcal{L}^{(5)} + \frac{1}{\Lambda^2} \mathcal{L}^{(6)} + O\left(\frac{1}{\Lambda^3}\right) , \quad (1)$$

where \mathcal{L}^{SM} is the usual SM lagrangian, $\mathcal{L}^{(5)}$ and $\mathcal{L}^{(6)}$ are the dimension five and six lagrangians and Λ is the energy scale for which one expects that whatever new physics exists beyond the SM becomes relevant. For our purposes - studies at the LHC - Λ will be of the order of one TeV. Note however that we will vary the coupling constants, denoted generically as $|a/\Lambda^2|$, from $10^{-12} \text{ TeV}^{-2}$ to 1 TeV^{-2} although some of the constraints from B physics will immediately discard some of values generated. This means that for an a of order 1, Λ is allowed to vary between 1 TeV and 10^6 TeV - 1 TeV just sets the scale of energy. The contributions from $\mathcal{L}^{(5)}$ break baryon and lepton number, and do not contribute to the FCNC physics we are interested in studying. The number of operators contained in $\mathcal{L}^{(6)}$ is enormous. Due to the gauge structure of the SM, a given dimension 6 operator that has an impact on top interactions can also have a parallel effect on processes involving only bottom quarks. B physics [7] is obviously the main source of constraints. The analysis of [7] as well as the one in [8] (see also [9]) can be used to impose limits on the sizes of several of our anomalous couplings, due to experimental constraints from B physics like $b \rightarrow s\gamma$ and others. There is a hierarchy of constraints resulting from the underlying SM gauge structure. The set of operators where the gauge structure manifests more strongly is the one denoted by LL in [7] as these operators are built with only $SU(2)$ doublets. Operators RR , built with singlets alone, are obviously the least constrained as no relation exists between an R -top and a R -bottom. The set of constraints obtained in [7] shows that operators of the type LL are already constrained beyond the reach of the LHC. Using the LHC [10, 11, 12] predictions they could show that some of the constraints on the operators coming from B physics are already stronger than what is predicted that could be measured at the LHC. This is true for operators of type LL , while limits on LR and RL operators are close to what is expected to be measured at the LHC. Hence, we discard all LL operators and we will use all the constraints on the dimension-6 operators obtained in [7]. B factories and the Tevatron are still collecting data and therefore these constraints will be even stronger by the time the LHC starts to analyse data. Finally we should mention that there are specific models for which some, but not all, effects on B physics of operators of type LL built with doublets only, cancel exactly [13]. In the model presented in that paper, anomalous contributions to vertices involving the Z boson and b quarks cancel precisely. However, the same does not occur for vertices with the W boson and b quarks - if we wish to cancel those contributions, we need to adjust the values of some of the anomalous couplings to make it happen. We still choose to leave those operators out of our analysis, but their inclusion may be of interest in the study of very specific models. Finally we note that B physics only constraints operators with origin in the electroweak sector.

The criteria mentioned above are very strict, and the number of operators which survive the selection

is small. For the strong sector, we have the operators [3]

$$\mathcal{O}_{itG} = i \frac{\alpha_{it}^S}{\Lambda^2} (\bar{u}_R^i \lambda^a \gamma_\mu D_\nu t_R) G^{a\mu\nu} \quad , \quad \mathcal{O}_{itG\phi} = \frac{\beta_{it}^S}{\Lambda^2} (\bar{q}_L^i \lambda^a \sigma^{\mu\nu} t_R) \tilde{\phi} G_{\mu\nu}^a \quad . \quad (2)$$

The coefficients α_{it}^S and β_{it}^S are complex dimensionless couplings, u_R^i and q_L^i represent the right-handed up-type quark and left-handed quark doublet of the first and second generation and $G_{\mu\nu}^a$ is the gluonic field tensor. There are also operators, with couplings α_{ti}^S and β_{ti}^S , where the positions of the top and u , q^i spinors are exchanged. The lagrangian obviously also includes the hermitian conjugates of these operators.

These operators generate FCNC vertices of the form $t q g$ or $t q g g$, among others (where $q = u, c$). The constants $\{\alpha^S, \beta^S\}$ are *a priori* unknown, the magnitude of which determines the importance of the FCNC processes. The Feynman rules for the FCNC vertices involving gluons stemming from these operators may be found in [3], and with those it is simple to calculate the decay width of $t \rightarrow q g$, given by (all quark masses except for the top's were set equal to zero)

$$\Gamma(t \rightarrow qg) = \frac{m_t^3}{12\pi\Lambda^4} \left\{ m_t^2 |\alpha_{tq}^S + (\alpha_{qt}^S)^*|^2 + 16 v^2 (|\beta_{tq}^S|^2 + |\beta_{qt}^S|^2) + 8 v m_t \text{Im} [(\alpha_{qt}^S + (\alpha_{tq}^S)^*) \beta_{tq}^S] \right\} \quad . \quad (3)$$

The electroweak FCNC interactions of the top quark were studied in [4], where it was shown that, according to the selection criteria described above, the operators describing them were given by

$$\begin{aligned} \mathcal{O}_{tB} &= i \frac{\alpha_{it}^B}{\Lambda^2} (\bar{u}_R^i \gamma_\mu D_\nu t_R) B^{\mu\nu} \quad , \quad \mathcal{O}_{tB\phi} = \frac{\beta_{it}^B}{\Lambda^2} (\bar{q}_L^i \sigma^{\mu\nu} t_R) \tilde{\phi} B_{\mu\nu} \quad , \\ \mathcal{O}_{tW\phi} &= \frac{\beta_{it}^W}{\Lambda^2} (\bar{q}_L^i \tau_I \sigma^{\mu\nu} t_R) \tilde{\phi} W_{\mu\nu}^I \quad , \quad \mathcal{O}_{\phi t} = \theta_{it} (\phi^\dagger D_\mu \phi) (\bar{u}_R^i \gamma^\mu t_R) \quad , \\ \mathcal{O}_{D_t} &= \frac{\eta_{it}}{\Lambda^2} (\bar{q}_L^i D^\mu t_R) D_\mu \tilde{\phi} \quad , \quad \mathcal{O}_{\bar{D}_t} = \frac{\bar{\eta}_{it}}{\Lambda^2} (D^\mu \bar{q}_L^i t_R) D_\mu \tilde{\phi} \quad . \end{aligned} \quad (4)$$

Once more, the complex constants $\{\alpha, \beta, \theta, \eta, \bar{\eta}\}$ are unknown, and used to parameterize the magnitude of the FCNC interactions.

One can use the equations of motion of the fields to establish relations between several of the operators shown above (see, for instance, eqs. (8) in [3] and more recently, the work of [14]). With our conventions, they translate, for the strong sector operators, as

$$\begin{aligned} \mathcal{O}_{utG}^\dagger &= -\mathcal{O}_{tuG} + \frac{i}{2} (\Gamma_u^\dagger \mathcal{O}_{utG\phi}^\dagger - \Gamma_t \mathcal{O}_{tuG\phi}) \\ \mathcal{O}_{utG}^\dagger &= \mathcal{O}_{tuG} + i g_s \bar{t} \gamma_\mu \gamma_R \lambda^a u \sum_i (\bar{u}^i \gamma^\mu \gamma_R \lambda_a u^i + \bar{d}^i \gamma^\mu \gamma_R \lambda_a d^i) \quad . \end{aligned} \quad (5)$$

where Γ_i are Yukawa couplings and the conventions of [6] were used. The first of these equations establishes that one can safely set to zero one of the anomalous couplings $\{\alpha, \beta\}$. The second one specifies that, if one were to include four-fermion operators in our calculation, one could set one further $\{\alpha, \beta\}$ coupling to zero - or, in exchange, to set one of the four-fermion couplings to zero. This is precisely what we have done in [4]. If we were analysing processes which only involved a top quark alongside gauge bosons (for instance, $t + Z$ production) it is arguable that one might then use both equations (5) and keep only two strong anomalous couplings, since no four-fermion terms contribute to those processes. However, in the work of refs. [3, 5], as well as in the current paper, we have performed a global analysis of the impact of FCNC top interactions. We have studied several channels simultaneously, namely top plus gauge boson production as well as four fermion channels (such as $q \bar{q} \rightarrow t \bar{q}$). As such, unless we were to include all four-fermion operators in our calculations, we are not allowed to use the second equation (5) to eliminate any coupling.

The importance of the simultaneous analysis of all channels is shown in ref. [5], where we computed the total anomalous cross section for single top plus jet production at the Tevatron - and then proceeded to use their experimental results to constrain the values of the anomalous couplings and several branching

ratios. Further, as can be seen in the expressions presented in [5], there are clearly *three* independent functions multiplying the several combinations of $\{\alpha, \beta\}$ couplings. This indicates, as discussed above, that by means of using the first equation of (5) we could “drop” one of those couplings, but not two of them. Nevertheless, we keep all $\{\alpha, \beta\}$ couplings since no significant simplification would come out of eliminating just one $\{\alpha, \beta\}$ coupling (none of the anomalous vertices shown in [3, 5] would be discarded). Finally, we found that this small coupling redundancy is quite useful as a cross-check of our computations.

Due to the standard Weinberg rotation we define new constants directly associated with the photon and the Z, given by

$$\begin{aligned}\alpha^\gamma &= \cos \theta_W \alpha^B, & \beta^\gamma &= \sin \theta_W \beta^W + \cos \theta_W \beta^B, \\ \alpha^Z &= -\sin \theta_W \alpha^B, & \beta^Z &= \cos \theta_W \beta^W - \sin \theta_W \beta^B.\end{aligned}\quad (6)$$

In terms of these anomalous couplings the decay widths for the decays $t \rightarrow q \gamma$ and $t \rightarrow q Z$ are given by [4]

$$\begin{aligned}\Gamma(t \rightarrow q \gamma) &= \frac{m_t^3}{64\pi\Lambda^4} \left\{ m_t^2 |\alpha_{tq}^\gamma + (\alpha_{qt}^\gamma)^*|^2 + 16 v^2 (|\beta_{tq}^\gamma|^2 + |\beta_{qt}^\gamma|^2) + \right. \\ &\quad \left. 8 v m_t \operatorname{Im} [(\alpha_{qt}^\gamma + (\alpha_{tq}^\gamma)^*) \beta_{tq}^\gamma] \right\}\end{aligned}\quad (7)$$

and

$$\begin{aligned}\Gamma(t \rightarrow q Z) &= \frac{(m_t^2 - m_Z^2)^2}{32 m_t^3 \pi \Lambda^4} \left[K_1 |\alpha_{qt}^Z|^2 + K_2 |\alpha_{tq}^Z|^2 + K_3 (|\beta_{qt}^Z|^2 + |\beta_{tq}^Z|^2) + K_4 (|\eta_{qt}|^2 + |\bar{\eta}_{qt}|^2) \right. \\ &\quad + K_5 |\theta|^2 + K_6 \operatorname{Re} [\alpha_{qt}^Z \alpha_{tq}^Z] + K_7 \operatorname{Im} [\alpha_{qt}^Z \beta_{tq}^Z] \\ &\quad + K_8 \operatorname{Im} [\alpha_{tq}^{Z*} \beta_{tq}^Z] + K_9 \operatorname{Re} [\alpha_{qt}^Z \theta^*] + K_{10} \operatorname{Re} [\alpha_{tq}^Z \theta] \\ &\quad \left. + K_{11} \operatorname{Re} [\beta_{qt}^Z (\eta_{qt} - \bar{\eta}_{qt})^*] + K_{12} \operatorname{Im} [\beta_{tq}^Z \theta] + K_{13} \operatorname{Re} [\eta_{qt} \bar{\eta}_{qt}^*] \right],\end{aligned}\quad (8)$$

with coefficients K_i given by

$$\begin{aligned}K_1 &= \frac{1}{2} (m_t^4 + 4 m_t^2 m_Z^2 + m_Z^4) & K_2 &= \frac{1}{2} (m_t^2 - m_Z^2)^2 & K_3 &= 4 (2 m_t^2 + m_Z^2) v^2 \\ K_4 &= \frac{v^2}{4 m_Z^2} (m_t^2 - m_Z^2)^2 & K_5 &= \frac{v^4}{m_Z^2} (m_t^2 + 2 m_Z^2) & K_6 &= (m_t^2 - m_Z^2) (m_t^2 + m_Z^2) \\ K_7 &= 4 m_t (m_t^2 + 2 m_Z^2) v & K_8 &= 4 m_t (m_t^2 - m_Z^2) v & K_9 &= -2 (2 m_t^2 + m_Z^2) v^2 \\ K_{10} &= -2 (m_t^2 - m_Z^2) v^2 & K_{11} &= -K_{10} & K_{12} &= -12 m_t v^3 & K_{13} &= \frac{-v^2}{m_Z^2} K_2.\end{aligned}\quad (9)$$

In expressions (3), (7) and (8), the couplings α_{qt}^S , etc, correspond to the separate cases $q = u$ and $q = c$. That is, the FCNC anomalous couplings will, in general, have different values for the decays of the “u” or “c” quarks.

III. FCNC CONTRIBUTIONS TO $BR(t \rightarrow b W)$

As we observe from eqs. (6), the fact that the Buchmüller and Wyler formalism is automatically gauge invariant imposes severe constraints on the type of new physics predicted by the effective operators we are considering. Namely, those equations show us that FCNC processes involving the photon cannot be entirely disentangled from those involving the Z, since the same operators which contribute to decays of the form $t \rightarrow q \gamma$ also contribute to $t \rightarrow q Z$. A careful analysis of the operators in eq. (4) shows us that some of them actually do contribute to decays of the form $t \rightarrow d W$ and $t \rightarrow s W$ - namely, the operators $\mathcal{O}_{tW\phi}$, \mathcal{O}_{D_t} and $\mathcal{O}_{\bar{D}_t}$, from whose structure one can “extract” the correct vertices [27]. The constants α do not contribute to these decays, as they stem from operators which involve only right-handed quarks.

Neither does the operator \mathcal{O}_{ϕ_t} , which has no d spinors. However, one can consider another operator of the type of \mathcal{O}_{ϕ_t} , namely

$$\mathcal{O}_{\phi_{td}} = i \frac{\theta_{td}}{\Lambda^2} (\phi^\dagger \epsilon D_\mu \phi) (\bar{t}_R \gamma_\mu d_R) \quad . \quad (10)$$

Gathering these operators, then, we arrive at the contribution of the FCNC top quark effective operators on its charged decays. This translates as a Feynman rule for “anomalous” top W interactions, presented in fig. (1). In this figure we consider only the vertex $W^+ t \bar{d}$, there is an analogous vertex for $W^+ t \bar{s}$ involving the anomalous “c” couplings. One important issue needs to be considered: none of these

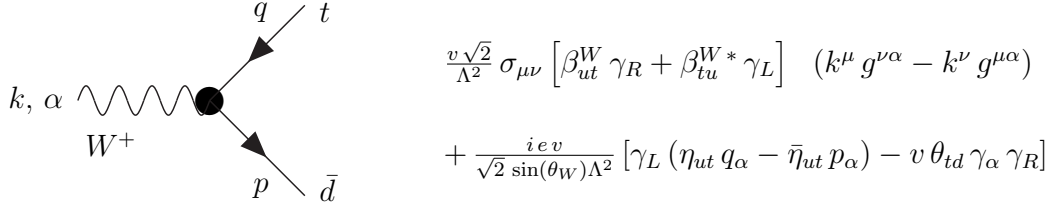


FIG. 1: Feynman rules corresponding to anomalous contributions to $t \rightarrow dW$.

anomalous operators has any contribution to the masses of the quarks. Thus, these contributions to the vertices $W t \bar{d}$ do not change the rotation between weak eigenstates of the quarks and their mass eigenstates. As such, these anomalous interactions constitute extra contributions to the CKM top matrix elements.

A subtlety now affects this calculation: in refs. [3, 4] we performed calculations of observables up to order $1/\Lambda^4$. Indeed, the FCNC decay widths of eqs. (3), (7) and (8) are of this order. However, that was due to the fact that, at tree-level, there are no Feynman diagrams contributing to the processes considered ($t \rightarrow qg$ decays, $t + Z$ production at the LHC, etc.). For the charged decays of the top, however, there are tree-level diagrams one must consider. Thus, the first and leading contribution to these processes stems from the interference terms between the SM process and the anomalous one, and is therefore of order $1/\Lambda^2$. No contributions of order $1/\Lambda^4$ will be considered here - these would have to include the square of the anomalous vertex of order $1/\Lambda^2$ (which we could compute easily) and the interference terms between the SM diagrams and those coming from effective operators of order $1/\Lambda^4$ (which we have not considered in this work).

Even though the anomalous vertex in fig. (1) includes contributions from several FCNC couplings, only β_{ut}^W (and β_{ct}^W) will have an impact on the decay $t \rightarrow dW$ (and $t \rightarrow sW$). This is due to the fact that we are considering all quark masses other than the top's equal to zero, and the chiral structure of the interactions in the vertex of fig. (1). A simple calculation yields,

$$\Gamma(t \rightarrow dW) = 1.42 |V_{td}|^2 - \frac{3g}{8\pi\Lambda^2} \frac{m_t^4 - m_W^4}{m_t^2} v \operatorname{Re}(V_{td} \beta_{ut}^W) \quad , \quad (11)$$

where the first term is the SM contribution [15], and the second one the interference with the anomalous operators. Thus, the branching ratio of $t \rightarrow bW$ will be given by

$$BR(t \rightarrow bW) = \frac{\Gamma(t \rightarrow bW)}{\sum_{q=u,c} [\Gamma(t \rightarrow qg) + \Gamma(t \rightarrow q\gamma) + \Gamma(t \rightarrow qZ)] + \sum_{q=d,s,b} \Gamma(t \rightarrow qW)} \quad , \quad (12)$$

where $\Gamma(t \rightarrow bW) = 1.42 |V_{tb}|^2 \text{ GeV}$ and the remaining widths are given by eqs. (3), (7), (8) and (11). This equation tells us that *all* the anomalous couplings considered so far have an impact on this branching ratio.

IV. ANALYSIS OF RESULTS

In refs. [3, 4] we have shown that the top quark anomalous FCNC interactions might have a large impact on single top production cross sections at the LHC. Furthermore, and due to the gauge-invariant nature of the formalism, we have shown that several FCNC observables are correlated. We also discovered a proportionality between several cross sections and certain FCNC widths. In this work we have shown that the same anomalous couplings which contribute to single top production at the LHC might also have a sizeable effect on a basic top quark property, its largest branching ratio. Thus, to probe the consistency of the formalism, an analysis of the expected deviations from the SM expected value, due to FCNC interactions, is in order.

A few words on the current state of knowledge about $Br(t \rightarrow bW)$: in the SM, at tree-level, this quantity is simply given by $|V_{tb}|^2$, assuming that the CKM matrix is unitary. Direct measurements of the top CKM coefficients are obviously quite difficult, and the corresponding results are affected by large error bars [17]. Nevertheless, the assumption of the unitarity of the CKM matrix is an extremely powerful one, and allows one to quote very precise values for these numbers which include all experimental results do date plus the theoretical input of unitarity. Namely, the PDG-listed values are [17]

$$V_{td} = 0.00874^{+0.00026}_{-0.00037}, \quad V_{ts} = 0.0407 \pm 0.0010, \quad V_{tb} = 0.999133^{+0.000044}_{-0.000043}. \quad (13)$$

As such, the SM prediction for $Br(t \rightarrow bW)$ is extremely precise, if one assumes the unitarity of the CKM matrix. All of our results which include FCNC interactions are therefore compared with the SM values under this assumption, and any eventual deviations will have to be interpreted in this light.

To investigate the importance that FCNC interactions might have on $Br(t \rightarrow bW)$, we made a vast scan of the possible values of the FCNC anomalous couplings. We considered almost two million different sets of anomalous couplings, varying them by several orders of magnitude, namely $10^{-12} \leq |a/\Lambda^2| \leq 1 \text{ TeV}^{-2}$, where a is a generic anomalous coupling. The upper limit in this variation comes from unitarity considerations [16]: given that the effective operator formalism is not renormalizable, it is not advisable to consider large values of the couplings, lest the growth of the anomalous contributions with \sqrt{s} becomes too steep. We also demanded that all partial FCNC anomalous branching fractions to be in agreement with all experimental bounds from LEP [18], Hera [19] and from the Tevatron [20, 21, 22]. For a detailed discussion about the referred experimental bounds see [5].

The same anomalous couplings which we use to compute FCNC branching ratios also enter into the expressions for FCNC processes of single top production, such as the $t + Z$ or $t + \text{jet}$ channels - those expressions may be found in refs. [3, 4, 5]. The cross sections were obtained using a cut on the p_T of the final state particles of 15 GeV. The partonic cross sections were integrated with CTEQ6M parton density functions [23], by setting the factorization scale equal to m_t , if the final state is a top quark and a jet, and $m_t + m_Z$, if the final state is $t + Z$. As we did in ref. [5], we will use the existing data from the Tevatron experiments concerning single top production to constrain our anomalous couplings. Namely, we will demand that, for a given set of anomalous couplings, the value of the extra FCNC contributions to the cross section for the single top + jet channel (which adds to its SM value) at the Tevatron be inferior to the observed experimental error for that observable [24], meaning 0.6 pb. This will eliminate a significant portion of our parameter space. We have also included all the constraints on the anomalous couplings derived from B physics in [7]. These constraints are really important as they further reduce the allowed parameter space improving the bounds on FCNC cross sections and branching ratios.

In figure 2 we plot the values of the FCNC cross section for $t + Z$ production at the LHC versus the values of $1 - Br(t \rightarrow bW)$. The central solid line corresponds to the SM-predicted value, assuming the CKM matrix is unitary. The error bar stemming from the values shown in eq. (13) is negligible. The points in red correspond to those combinations of anomalous couplings which were rejected due to the experimental constraints from the Tevatron, which for this observable do not exclude a significant region of the plot. The blue points are the ones excluded by the B physics analysis in [7]. The most important aspect of this plot is the fact that, if one were to measure at the LHC, with some precision, the value of the branching ratio for $t \rightarrow bW$ at the LHC and find a value very close to that which is to be expected if the SM holds, then one would expect the cross section for $t + Z$ production at the LHC to be smaller than about 1 pb. In fact, the expected SM value for the branching ratio provides us an upper bound on that cross section. Alternatively, if one were to measure a cross section for $t + Z$ production larger than approximately 1 pb, then the assumption that presides the calculation of $Br(t \rightarrow bW)$ in this figure would be incorrect - meaning, the values of the top CKM matrix elements would be quite different from those of eq. (13), and the CKM matrix could not be considered unitary.

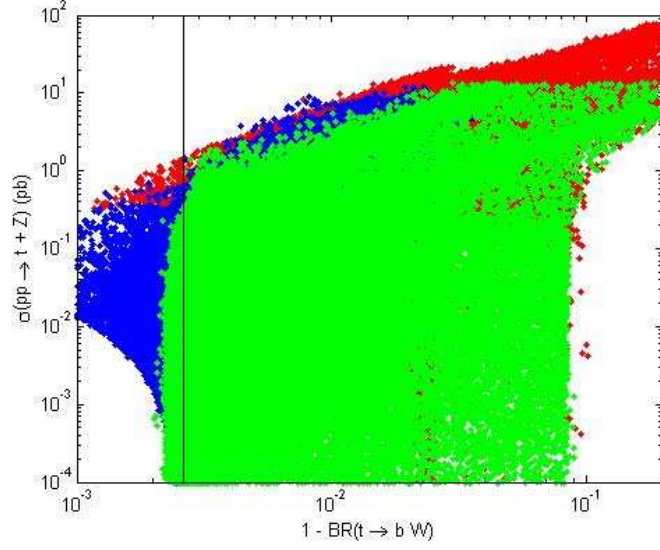


FIG. 2: FCNC cross section for $t + Z$ production at the LHC, versus $1 - Br(t \rightarrow bW)$. The red points correspond to those portions of parameter space excluded by the Tevatron results, while the blue points are excluded by the B physics analysis in [7]. The solid line is the SM-predicted value for $1 - Br(t \rightarrow bW)$ and the points in green are those which comply with both the Tevatron and B-factories data.

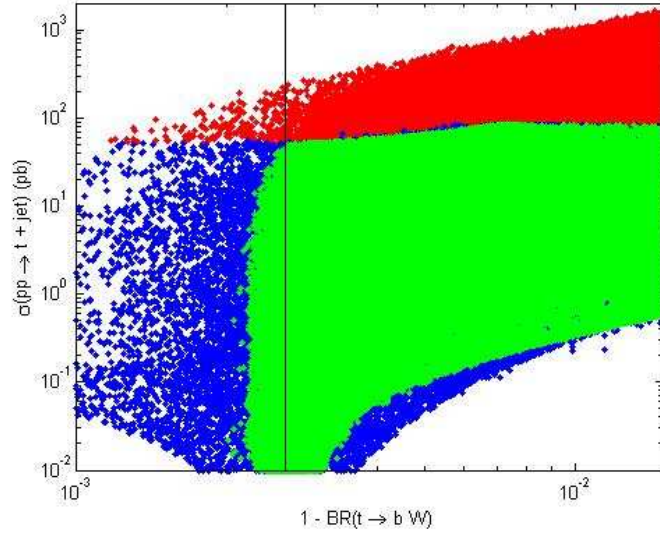


FIG. 3: FCNC contributions to cross section for $t + jet$ production at the LHC, versus $1 - Br(t \rightarrow bW)$. The red points correspond to those portions of parameter space excluded by the Tevatron results, while the blue points are excluded by the B physics analysis in [7]. The solid line is the SM-predicted value for $1 - Br(t \rightarrow bW)$ and the points in green are those which comply with both the Tevatron and B-factories data.

Consider now figure 3, where we plot the total FCNC contribution to the top + jet cross section [28] at the LHC. The points which survive the Tevatron and the B-factories exclude a significant portion of available parameter space. The line corresponding to the SM prediction for $1 - Br(t \rightarrow bW)$ once again establishes an upper bound for the FCNC contribution to this cross section, namely about 40 pb. The predicted SM value for single top production at the LHC is around 202 pb, with a theoretical error of 8 pb [25]. This includes the $t + W$ channel. If one considers only the s and t channel contributions (so as to obtain an estimate for the SM top + jet cross section) the value is $\sigma^{SM}(pp \rightarrow t + jet) = 168 \pm 6$ pb. In this case the t -channel

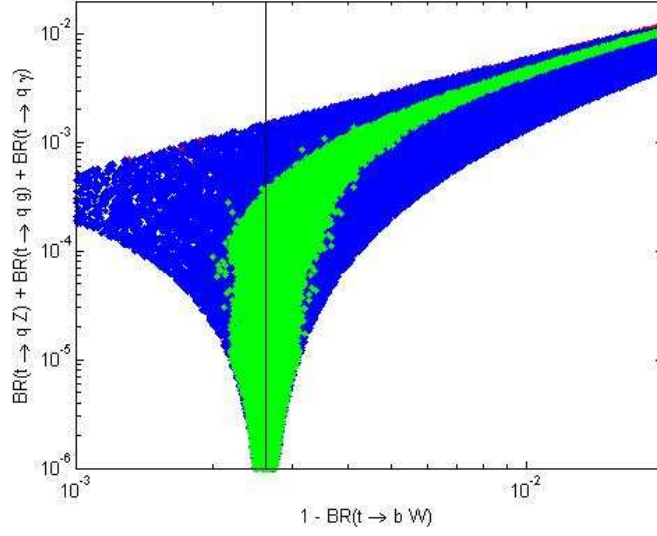


FIG. 4: Total FCNC branching ratios versus $1 - Br(t \rightarrow b W)$. The points in green are those which comply with both the Tevatron and B-factories data. The red points correspond to parameter space excluded by the Tevatron and the blue points are excluded by the B physics analysis in [7]

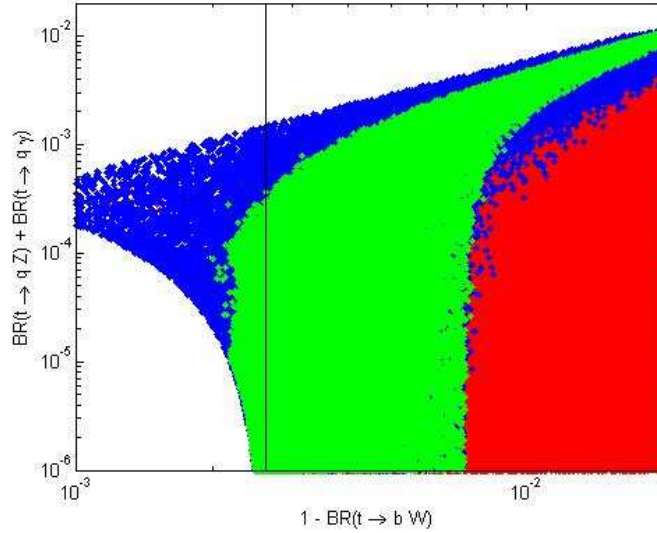


FIG. 5: Electroweak FCNC branching ratios versus $1 - Br(t \rightarrow b W)$. The red points correspond to those portions of parameter space excluded by the Tevatron results, while the blue points are excluded by the B physics analysis in [7]. The solid line is the SM-predicted value for $1 - Br(t \rightarrow b W)$ and the points in green are those which comply with both the Tevatron and B-factories data.

cross section is expected to be measurable with a total error of 10 % at the LHC while an uncertainty of 36 % was estimated for the s-channel mode [26]. Therefore, the error in this quantity will probably be inferior to the maximum FCNC contribution to it, taken from fig. 3. A not very large deviation from the SM value for this cross section would therefore imply the presence of sizeable FCNC contributions, even with a measurement of $Br(t \rightarrow b W)$ in full agreement with its expected SM values.

Finally, we may use the current values for $Br(t \rightarrow b W)$, assuming unitarity of the CKM matrix, to ask: what is the maximum value for the FCNC branching ratios that is allowed under those assumptions? The answer lies in fig. 4, where we plot the sum of all top FCNC branching ratios versus $1 - Br(t \rightarrow b W)$.

The data shown in this plot already takes the Tevatron results on single top production into account as well as the constraints from B-physics. Results from the Tevatron are very important but the results from B-physics are even more restrictive and this is in full display in fig. 5. In this figure we plotted $1 - Br(t \rightarrow bW)$ against the electroweak FCNC branching ratios - we can appreciate just how much parameter space is excluded by the B-physics results from [7].

Reading the intersection between the SM-predicted value for $1 - Br(t \rightarrow bW)$ and the band of values allowed for the FCNC branching ratios from fig. 4, we obtain the following upper bound for the total FCNC branching ratio:

$$Br(t \rightarrow qX) < 4.0 \times 10^{-4} \quad (14)$$

while using only Tevatron data the constraint would be 1.0×10^{-3} , again showing the relevance of the B-physics results. If we now read the intersection between the SM-predicted value for $1 - Br(t \rightarrow bW)$ and the band of values allowed for the FCNC branching ratios from fig. 5, we obtain a very similar upper bound for the electroweak FCNC branching ratio:

$$Br(t \rightarrow qZ) + Br(t \rightarrow q\gamma) < 4.0 \times 10^{-4} \quad . \quad (15)$$

The best current upper bound for the electroweak sector are the Hera and LEP for $Br(t \rightarrow q\gamma)$ and from the Tevatron for $Br(t \rightarrow qZ)$ and are all at the percent level. Therefore our results are better by at least one order of magnitude. Moreover they are obtained for the sum of branching ratios with $q = u, c$. Therefore they apply to each individual branching ratio.

The excellent results obtained by the Tevatron from direct top production [20] agree with ours for $q = u$ and our results are slightly better for $q = c$. We remind the reader that our results were obtained under one important theoretical assumption namely, the unitarity of the CKM matrix.

V. CONCLUSIONS

We considered a vast set of dimension six effective operators contributing to FCNC interactions of the top quark. In previous works the effects of those operators in processes of single top production was analysed in great detail. In this letter, we concerned ourselves with a more basic observable, the main branching ratio of the top, $Br(t \rightarrow bW)$. Several of the operators that contribute to FCNC decays of the top also affect its charged decays. We thus have contributions from FCNC anomalous couplings to top charged branching ratios, through interference terms between SM Feynman diagrams and diagrams containing an anomalous vertex. Those FCNC contributions would then constitute corrections to the top quark CKM matrix elements.

We computed $Br(t \rightarrow bW)$ for a vast set of possible values for the FCNC anomalous couplings, and analysed the cross sections for several processes of single top production via FCNC at the LHC. The values predicted by the SM for $Br(t \rightarrow bW)$, if one assumes unitarity of the CKM matrix, are very precise and give us an immediate upper bound for the FCNC cross sections. A measurement of a LHC cross section for $t + Z$ production at the LHC superior to about 1 pb, for instance, would imply that the CKM matrix is in fact non-unitary.

Assuming unitarity of the CKM matrix also allows us to obtain a rather good upper bound on the total FCNC branching ratio, better than most current direct experimental measurements for operators in the electroweak sector. For operators stemming from the strong sector, the CDF results are competitive with ours especially the ones involving the u -quark. The curious interplay between top FCNC interactions and its charged decays, presented in this paper, allows a consistency check between several top quarks observables. If the top quark indeed has sizeable FCNC interactions, stemming from some type of new physics, the correlations between these several observables will provide us with a better understanding of those interactions.

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- [1] J. A. Aguilar-Saavedra, *Acta Phys. Polon. B* **35** (2004) 2695 [arXiv:hep-ph/0409342].
 - [2] M.E. Luke and M.J. Savage, *Phys. Lett.* **B307** (1993) 387;
D. Atwood, L. Reina and A. Soni, *Phys. Rev.* **D55** (1997) 3156;
J.M. Yang, B.L. Young and X. Zhang, *Phys. Rev.* **D58** (1998) 055001;
J. Guasch and J. Solà, *Nucl. Phys.* **B562** (1999) 3;
D. Delepine and S. Khalil, *Phys. Lett.* **B599** (2004) 62;
J.J. Liu, C.S. Li, L.L. Yang and L.G. Jin, *Phys. Lett.* **B599** (2004) 92;
G. Eilam, M. Frank and I. Turan, *Phys. Rev.* **D74** (2006) 035012;
J. J. Cao, G. Eilam, M. Frank, K. Hikasa, G. L. Liu, I. Turan and J. M. Yang, *Phys. Rev.* **D75** (2007) 075021;
A. Arhrib and W. S. Hou, *JHEP* **0607** (2006) 009;
A. Arhrib, K. Cheung, C. W. Chiang and T. C. Yuan, *Phys. Rev.* **D73** (2006) 075015;
J. Guasch, W. Hollik, S. Penaranda and J. Sola, *Nucl. Phys. Proc. Suppl.* **157** (2006) 152;
D. Lopez-Val, J. Guasch and J. Sola, hep-ph/0710.0587 ;
J. A. Aguilar-Saavedra, *Phys. Rev.* **D67** (2003) 035003 [Erratum-ibid. **D69** (2004) 099901];
F. del Aguila, J. A. Aguilar-Saavedra and R. Miquel, *Phys. Rev. Lett.* **82** (1999) 1628;
T. P. Cheng and M. Sher, *Phys. Rev.* **D35** (1987) 3484;
S. Bejar, J. Guasch and J. Sola, *Nucl. Phys.* **B600** (2001) 21;
C. S. Li, R. J. Oakes and J. M. Yang, *Phys. Rev.* **D49** (1994) 293 [Erratum-ibid. **D56** (1997) 3156];
G. M. de Divitiis, R. Petronzio and L. Silvestrini, *Nucl. Phys.* **B504** (1997) 45
J. L. Lopez, D. V. Nanopoulos and R. Rangarajan, *Phys. Rev.* **D56** (1997) 3100;
G. Eilam, A. Gemintern, T. Han, J. M. Yang and X. Zhang, *Phys. Lett.* **B510** (2001) 227.
 - [3] P. M. Ferreira, O. Oliveira and R. Santos, *Phys. Rev.* **D73** (2006) 034011;
P. M. Ferreira and R. Santos, *Phys. Rev.* **D73** (2006) 054025;
P. M. Ferreira and R. Santos, *Phys. Rev.* **D74** (2006) 014006;
 - [4] P. M. Ferreira, R. B. Guedes and R. Santos, *Phys. Rev.* **D77** (2008) 114008.
 - [5] R. A. Coimbra, P. M. Ferreira, R. B. Guedes, O. Oliveira, A. Onofre, R. Santos and M. Won, *Phys. Rev. D* **79** (2009) 014006 [arXiv:0811.1743 [hep-ph]].
 - [6] W. Buchmüller and D. Wyler, *Nucl. Phys.* **B268** (1986) 621.
 - [7] P. J. Fox, Z. Ligeti, M. Papucci, G. Perez and M. D. Schwartz, arXiv:0704.1482 [hep-ph].
 - [8] B. Grzadkowski and M. Misiak, *Phys. Rev. D* **78** (2008) 077501 [arXiv:0802.1413 [hep-ph]].
 - [9] Y. Grossman, Z. Ligeti and Y. Nir, arXiv:0904.4262 [hep-ph].
 - [10] J. Carvalho et al., *Eur. Phys. J.C* **52** (2007) 999-1019.
 - [11] T. Lari et al., Report of Working Group 1 of the CERN Workshop “Flavour in the era of the LHC”, hep-ph/0801.1800.
 - [12] CMS Physics TDR: Volume II, CERN/LHCC 2006-021, <http://cmsdoc.cern.ch/cms/cpt/tdr/>.
 - [13] F. del Aguila, M. Pérez-Victoria, J. Santiago, *Phys. Lett.* **B492** (2000) 98; *id.*, *JHEP* **09** (2000) 011.
 - [14] J. A. Aguilar-Saavedra, hep-ph/0811.3842.
 - [15] A. Denner and T. Sack, *Nucl. Phys.* **B358** (1991) 46;
G. Eilam, R.R. Mendel, R. Migneron and A. Soni, *Phys. Rev. Lett.* **66** (1991) 3105;
A. Czarnecki and K. Melnikov, *Nucl. Phys.* **B554** (1999) 520;
K.G. Chetyrkin, R. Harlander, T. Seidensticker and M. Steinhauser, *Phys. Rev.* **D60** (1999) 114015;
S.M. Oliveira, L. Brucher, R. Santos and A. Barroso, *Phys. Rev.* **D64** (2001) 017301.
 - [16] J. Wudka, *Int. J. Mod. Phys. A* **9** (1994) 2301.
 - [17] C. Amsler et al, *Phys. Lett.* **B667** (2008) 1.
 - [18] ALEPH Coll., A. Heister et al., *Phys. Lett.* **B543** (2002) 173;
DELPHI Coll. J. Abdallah et al. *Phys. Lett.* **B 590** (2004) 21;
OPAL Coll., G. Abbiendi et al., *Phys. Lett.* **B 521** (2001) 181;
L3 Coll., P. Achard et al., *Phys. Lett.* **B 549** (2002) 290; M. Beneke et al., “Top quark physics”, in “Standard Model physics (and more) at the LHC”, G. Altarelli and M. L. Mangano eds., Geneva, Switzerland: CERN (2000), [arXiv:hep-ph/0003033].
 - [19] ZEUS Coll., S. Chekanov et al., *Phys. Lett.* **B 559** (2003) 153; A. A. Ashimova and S. R. Slabospitsky, H1 Coll., A. Aktas et al., *Eur. Phys. J.C* **33**, (2004), 9.
 - [20] T. Aaltonen et al. [CDF Collaboration], *Phys. Rev. Lett.* **102** (2009) 151801 [arXiv:0812.3400 [hep-ex]].
 - [21] T. Aaltonen et al. [CDF Collaboration], *Phys. Rev. Lett.* **101** (2008) 192002 [arXiv:0805.2109 [hep-ex]]; <http://www-cdf.fnal.gov/physics/new/top/2008/tprop/TopFCNC-v1.5/>
 - [22] V. M. Abazov et al. [D0 Collaboration], *Phys. Rev. Lett.* **99** (2007) 191802; CDF Coll., F. Abe et al., *Phys. Rev. Lett.* **80** (1998) 2525; CDF Coll., V. M. Abazov et al. *Phys. Rev. Lett.* **99**, 191802 (2007).
 - [23] J. Pumplin et al, *JHEP* **0207** (2002) 012.
 - [24] T. Aaltonen et al. [The CDF collaboration], arXiv:0903.0885 [hep-ex]; V. M. Abazov et al. [The D0 Collab-

- oration], arXiv:0903.0850 [hep-ex].
- [25] W. Bernreuther, J. Phys. G **35** (2008) 083001 and references therein;
N. Kidonakis, Phys. Rev. D **74** (2006) 114012;
N. Kidonakis, Phys. Rev. D **75** (2007) 071501.
 - [26] G. L. Bayatian et al. [CMS Collaboration], J. Phys. G 34, 995 (2007); ATLAS detector and physics performance. Technical design report. Vol. 2,, report CERN-LHCC-99-15.
 - [27] The absence of any anomalous contributions to the decay $t \rightarrow bW$ is a direct consequence of the physical criteria we applied to the selection of the effective operators. Namely, the exclusion of any operators having a direct impact on bottom quark physics.
 - [28] As was explained in refs. [3, 4, 5], the FCNC top + jet cross section contributions add to those expected from the SM.